

Radiatively decaying scalar dark matter through U(1) mixings and the Fermi 130 GeV gamma-ray line

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Abstract

In light of the recent observation of the Fermi-LAT 130 GeV gamma-ray line, we suggest a model of scalar dark matter in hidden sector, which can decay into two (hidden) photons. The process is radiatively induced by a GUT scale fermion in the loop, which is charged under a hidden sector U(1), and the kinetic mixing ($\sim \epsilon F^{\mu\nu} F'_{\mu\nu}$) enables us to fit the required decay width for the Fermi-LAT peak. The model does not allow any dangerous decay channels into light standard model particles.

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I. INTRODUCTION

Although the major constituent of matter content of the Universe is dark matter (DM), we only know little about its nature so far [1]. Having no DM candidate in its particle contents, the standard model (SM) is strongly required to be extended. One intriguing possibility is that a hidden sector attached to the standard model sector includes a dark matter candidate. The singlet DM candidate could be a scalar [2], a fermion [3–5] or a vector boson [6, 7].

The recent claim of the Fermi Large Area Telescope (Fermi-LAT) [8] 130 GeV γ -ray line [9, 10] shed light on further details of the dark matter property since no known astrophysical source would produce such a peak. The claim was further strengthened by [11] (and also [12, 13]) even though the Fermi-LAT collaboration only has provided sensitivity limits on dark matter models based on a part of the acquired data set on different region of interest [14]. There are explanations of this γ -ray line based on spectral and spatial variations of diffuse γ -ray [15] and new background with ‘Fermi-bubble’ [16], but the most interesting interpretation might be that the γ -ray line could be originated from the DM annihilation, $\chi\chi \rightarrow \gamma\gamma$ (or γX) with $E_\gamma \simeq m_\chi$ (or $m_\chi(1 - m_X^2/4m_\chi^2)$) $\simeq 130$ GeV [17–29] or the DM decay, $\chi \rightarrow \gamma\gamma$ (or γX) with $2E_\gamma \simeq m_\chi$ (or $m_\chi(1 - m_X^2/m_\chi^2)$) $\simeq 260$ GeV [27, 30].¹

For the DM annihilation interpretation of the 130 GeV γ -ray peak, the required values for annihilation cross section is found to be $\langle\sigma v\rangle_{\chi\chi\rightarrow\gamma\gamma} \sim \text{a few} \times 10^{-27} \text{ cm}^3/\text{s}$,² which is approximately one order of magnitude smaller than the total annihilation cross section for the thermal production of DM, $\langle\sigma v\rangle_{\chi\chi\rightarrow SM} \simeq 3 \times 10^{-26} \text{ cm}^3/\text{s}$ [1]. As the dark matter is likely to be electrically neutral (or milli-charged [4, 5, 34]), the annihilation process for $\gamma\gamma$ production may be radiatively induced by massive charged particles in the loop. If some of the charged particles are lighter than the DM particle, there could appear tree-level annihilation channels to these charged particles, which may dominantly determine the relic abundance of dark matter. However, the loop factor is too small as $g^2/16\pi^2 \lesssim 10^{-2}$ and thus does not correctly account the discrepancy between the cross sections. A variety of

¹ In Refs. [31–33], it has been studied to constrain the DM models accompanied by continuum photons correlated to the 130 GeV gamma-ray line.

² More precisely, $\langle\sigma v\rangle_{\chi\chi\rightarrow\gamma\gamma} \simeq (1.27 \pm 0.32_{-0.28}^{+0.18}) \times 10^{-27} \text{ cm}^3/\text{s}$ ($2.27 \pm 0.57_{-0.51}^{+0.32} \times 10^{-27} \text{ cm}^3/\text{s}$) for the Einasto (NFW) DM profile [10].

annihilating DM models have been suggested to overcome this issue [17–29].

Decaying DM can be an alternative explanation. Indeed, decaying dark matter models have been recently proposed to account the excessive observation of positron in the PAMELA and ATIC where the dark matter is a vector boson in a hidden sector [7]. The vector boson of the hidden sector Abelian gauge group $U(1)'$ can decay to the standard model photon through the kinematical mixing term $\epsilon F_{\mu\nu} F'^{\mu\nu}$, where $F^{\mu\nu}$ ($F'^{\mu\nu}$) is the field strength tensor of $U(1)$ ($U(1)'$) gauge boson, respectively. As the mixing parameter could be small ($\epsilon \sim 10^{-26}$) [7], the decay width could be suppressed. However, the decay of vector boson to a pair of photons is forbidden by the Landau-Yang theorem [35] so that we need another model. In Refs. [27, 30], a scalar dark matter, ϕ , was considered with an effective operator allowing the decay to two photons: $c_6 \frac{\phi}{\Lambda^2} F^{\mu\nu} F_{\mu\nu}$, which is dimension six.³ It is pointed out in Ref. [30] that a dimension five operator, $c_5 \frac{\phi}{\Lambda} F^{\mu\nu} F_{\mu\nu}$, cannot fit the data without introducing Trans-Planckian cutoff ($\Lambda \gg M_{Pl}$) or equivalently a largely suppressed coefficient $c_6 \ll 1$ as the required partial decay width of the dark matter to photons is extremely small, $\Gamma(\phi \rightarrow \gamma\gamma) \sim 10^{-29} s^{-1}$.

In this paper, we try to combine the advantages of above two cases:

- A scalar dark matter *can* decay into $\gamma\gamma$ differently from the massive vector dark matter,
- A small kinetic mixing ϵ can make the effective couplings of the dark matter particle with the standard model particles small.

Combining these two advantages, we suggest a dark matter model, which has an effective dimension five operators of the form:

$$\mathcal{O} = c_5 \frac{\phi}{\Lambda} (F'_{\mu\nu} F'^{\mu\nu} + \epsilon F_{\mu\nu} F'^{\mu\nu} + \epsilon^2 F_{\mu\nu} F^{\mu\nu}), \quad (1)$$

from which we can learn that the decay amplitudes to $\gamma\gamma'$ and $\gamma\gamma$ are relatively suppressed by a factor of ϵ and ϵ^2 with respect to the one for $\gamma'\gamma'$ channel due to the mixing.

In the next section (Sec. II), we further explain the model in detail and present the partial decay widths of the dark matter to (hidden) photons then clarify the model parameter space providing a good fit to the 130 GeV gamma-ray line. Discussions on possible experimental bounds on the same parameter space follow. In Sec. III, we further discuss the theoretical

³ For radiative DM decays to $\gamma\gamma$ or γX , see e.g. [36].

issues concerning the consistency of the model and also other cosmological observations then conclude in Sec. IV. Finally, in Appendix, we present some details of the U(1) mixing Lagrangian for an unbroken U(1) symmetry.

II. THE MODEL AND EXPERIMENTAL BOUNDS

Postulating extra U(1) gauge symmetries is one of the simplest extensions of the standard model. As the kinetic mixing term $\epsilon F^{\mu\nu} F'_{\mu\nu}$ is compatible with Lorentz as well as gauge symmetry, the term should be included in view of effective field theory. The term can be generated through one-loop diagrams with a bi-charged fermion [37]. If the extra U(1) is broken by a hidden sector Higgs mechanism, the gauge boson (Z') gets mass and mixes with the standard model Z -boson [38]. A general analysis for a hidden sector DM, which is charged under a broken $U(1)_H$ was done in Ref. [39]. On the other hand, if the extra U(1) is exact, the massless gauge boson (hidden photon or exphoton, γ_H) can mix with the usual photon (γ) and the corresponding phenomenology becomes quite different from the case with Z' . It is worth noticing that a light DM explanation with the hidden photon for anomalous 511 keV γ -ray signature from the Galactic Center was considered in Ref. [5], where a milli-charged fermion dark matter was introduced.

For the Fermi-LAT γ -ray peak, a scalar DM is more profitable. The minimal setup only introduces a heavy vector-like fermion ψ , which is only charged under the $U(1)_H$, and a scalar dark matter candidate ϕ . The hidden sector can interact with the SM sector through a kinetic mixing between $U(1)_{EM}$ and $U(1)_H$.⁴ The model Lagrangian is given by

$$\mathcal{L} \supset \mathcal{L}_{SM} - \frac{1}{4} \hat{F}_{H\mu\nu} \hat{F}_H^{\mu\nu} - \frac{\sin \epsilon}{2} \hat{F}_{\mu\nu} \hat{F}_H^{\mu\nu} - \lambda \phi \bar{\psi} \psi + i \bar{\psi} \gamma^\mu (\partial_\mu - i \hat{g}_H \hat{A}_\mu^H) \psi - m_\psi \bar{\psi} \psi, \quad (2)$$

where $\hat{F}^{\mu\nu}$ and $\hat{F}_H^{\mu\nu}$ are respectively field strength tensors for $U(1)_{EM}$ and $U(1)_H$. The detailed derivation of the kinetic mixing between $U(1)_{EM}$ and massless $U(1)_H$ can be found in Appendix A. Current bounds on hidden U(1) gauge bosons and millicharged particles (MCPs) are well summarized in Ref. [40]. Indeed, when $m_{\gamma_H} = 0$, the only change for the SM photon interactions is the modification of the coupling constant, $\hat{e} \rightarrow \hat{e}/\cos \epsilon$, which can be simply refined by field redefinition.

⁴ Here the ‘Higgs-portal’ interaction, $\phi^2 |H|^2$, is assumed to be negligible as we do not allow the DM decay via mixing with the Higgs boson. Further discussion is given in Sec. III.

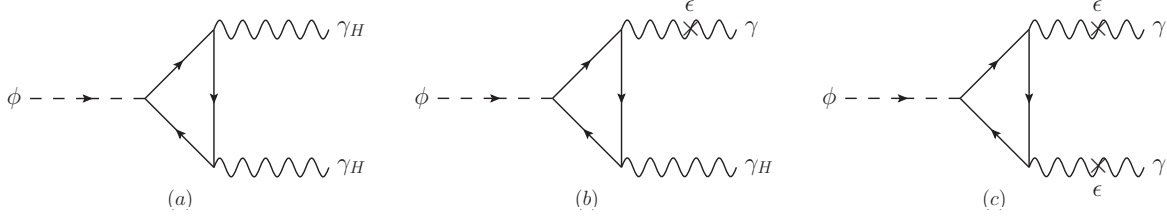


FIG. 1: (a) Scalar dark matter decaying to two hidden photons. (b) Scalar dark matter decaying to one hidden photon and one SM photon. (c) Scalar dark matter decaying to two SM photons.

The scalar DM candidate ϕ can radiatively decay to hidden photon ($A_H = \gamma_H$) as well as the conventional photon (γ) of the EM interaction through the triangle diagrams with the virtual hidden sector fermion ψ as can be seen from Fig. 1. From the couplings in Eq. (2), one can easily calculate the decay width of ϕ :

$$\Gamma(\phi \rightarrow \gamma_H \gamma_H) = \frac{(\alpha_H \lambda)^2}{256\pi^3} \frac{m_\phi^3}{m_\psi^2} |F(\tau)|^2, \quad (3)$$

where $\alpha_H = g_H^2/4\pi$ and $F(\tau = 4m_\psi^2/m_\phi^2) = -2\tau [1 + (1 - \tau) \arcsin^2(1/\sqrt{\tau})]$ which is well approximated by $-4/3$ at a large τ limit. To make ϕ stable enough, we would require the longevity of ϕ :

$$\Gamma(\phi \rightarrow \gamma_H \gamma_H)^{-1} \gg \tau_{\text{Universe}} \approx 4.34 \times 10^{17} \text{ sec}. \quad (4)$$

Thus, we can obtain a constraint on the combination $\alpha_H \lambda$:

$$\alpha_H \lambda \ll 1.96 \times 10^{-7} \left(\frac{m_\psi}{10^{16} \text{ GeV}} \right) \quad \text{for} \quad m_\phi = 260 \text{ GeV}. \quad (5)$$

The DM candidate ϕ predominantly decays to two hidden photons, $\phi \rightarrow \gamma_H \gamma_H$ [Fig. 1-(a)], which thus determines the life time of ϕ but a hidden photon can be converted into a SM photon through the kinetic mixing. Consequently, one can detect the DM decay signals through the decay modes $\phi \rightarrow \gamma_H \gamma / \gamma \gamma$ [Fig. 1-(b)/(c)], which are respectively suppressed by ϵ^2 and ϵ^4 compared with the dominant decay mode:

$$\Gamma(\phi \rightarrow \gamma_H \gamma_H) : \Gamma(\phi \rightarrow \gamma_H \gamma) : \Gamma(\phi \rightarrow \gamma \gamma) \simeq 1 : \epsilon^2 : \epsilon^4. \quad (6)$$

This helps as the effective dimension five operators, $\phi F_{\mu\nu} F^{\mu\nu}$ and $\phi F_{\mu\nu} X^{\mu\nu}$, are all suppressed by powers of ϵ and provides the required decay rate $\Gamma^{-1} \approx \mathcal{C} \times 10^{29} \text{ sec}$ [30, 31]

or

$$\Gamma(\phi \rightarrow \gamma_H \gamma)^{-1} \approx 1.52\mathcal{C} \times 10^{53} \text{ GeV}^{-1}, \quad (7)$$

where a convenient parameter $\mathcal{C} \in (0.1, 1)$ is introduced. Then, the parameter range for the mixing parameter ϵ can be read:

$$\epsilon \approx \frac{4.1 \times 10^{-13}}{\alpha_H \lambda \sqrt{\mathcal{C}}} \left(\frac{m_\psi}{10^{16} \text{ GeV}} \right) \quad \text{for } m_\phi = 260 \text{ GeV}. \quad (8)$$

Having a small value for $\alpha_H \lambda \ll 10^{-7}$, a relatively sizable kinetic mixing parameter ϵ is required.

In Fig. 2, we plotted the parameter space in the $m_\psi - \alpha_H \lambda$ plane. The upper left region in shade (grey) is excluded by the longevity of the dark matter. The three colored bands are for fitting the required decay width ($\Gamma^{-1} = 10^{28} - 10^{29} \text{ sec}$) with $\epsilon = 10^{-4}, 10^{-3}$ and 10^{-2} , respectively. The charged fermion ψ is assumed to be heavy ($\sim 10^{16} \text{ GeV}$).

III. OTHER ISSUES

In this section, we discuss possible difficulty of the ‘Higgs-portal’ type interaction and its way out. Then, we propose some scenarios for producing the singlet dark matter from the inflaton or other heavy particle decays in the early universe.

A. $\sigma\phi^2|H|^2$ or $\mu\phi|H|^2$ type interactions

The gauge invariant renormalizable interactions between the dark matter and the Higgs boson of the form $\sim \sigma\phi^2|H|^2$ (“Higgs portal”) and $\mu\phi|H|^2$ are generically allowed. If $m_\phi \geq 2m_h$, ϕ can therefore decay into two Higgs bosons with the decay rate:

$$\Gamma(\phi \rightarrow hh) \propto \frac{1}{16\pi} \frac{(\sigma\langle\phi\rangle + \mu)^2}{m_\phi} \sqrt{1 - \left(\frac{2m_h}{m_\phi}\right)^2}, \quad (9)$$

which can be significantly larger than the required decay width in Eq. (4). This difficulty can be avoided when an additional spatial dimension exists and all the hidden sector particles (ϕ, ψ, γ_H) are localized on the ‘hidden-brane’ which is spatially distant from the ‘visible-brane’ on which all the SM particles ($\supset \{H, \gamma, \text{leptons}, \text{quarks}\}$) reside. A mediator fermion ψ_M , which is charged under the both of the Abelian gauge groups $U(1)_{\text{EM}} \times U(1)_H$, could

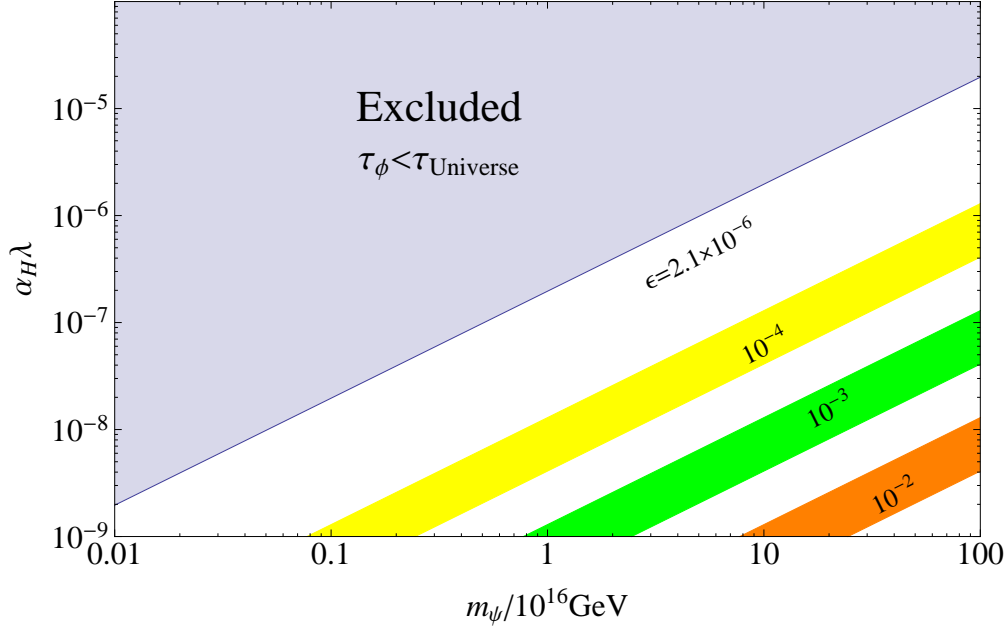


FIG. 2: Contour plot for the kinetic mixing parameter ϵ in the $m_\psi - \alpha_H \lambda$ plane. The colored region above the thick line is excluded since the life time of ϕ is shorter than the age of the Universe. The partial life time of ϕ to $\gamma\gamma_H$ is fixed to be in $(10^{28}, 10^{29})$ sec, which is required to fit the Fermi-LAT data, with $\epsilon = 10^{-4} - 10^{-2}$ in the colored bands from the top to the bottom on the right lower side of the graph. The larger $\alpha_H \lambda$ is required for the smaller ϵ .

provide a sizable $\gamma - \gamma_H$ mixing. Then, the operator $\hat{\mathcal{O}} \sim \phi^2 H H^\dagger$ is suppressed by many loop factors (two one-loops of the $\phi\phi - \gamma_H \gamma_H$ and $HH - \gamma\gamma$ vertices, and two additional one-loops of the $\gamma - \gamma_H$ mixings) and also by possible suppressed wave function overlaps $\sim \psi(y_{\text{hidden}})^* \psi(y_{\text{visible}})$. Similarly, the other operator $\hat{\mathcal{O}} \sim \phi H H^\dagger$ is also suppressed.

B. The production of ϕ

1. Model-1

A simple mechanism for the production of ψ and ϕ is the inflaton decay after reheating: $\Phi \rightarrow \psi\bar{\psi}$ and $\Phi \rightarrow \phi\phi$, where Φ is the inflaton field. As the mass of ψ is around the GUT scale, we do not expect any sizable number of remaining ψ in our patch universe, but still

ϕ could remain to be a good DM candidate. Here the only requirement is

$$m_\psi > T_{\text{reheating}} > m_\phi, \quad (10)$$

where a large space of reheating temperature in TeV to GUT-scale could be allowed, in principle. With this condition, ϕ can be produced but ψ cannot. In this paper, we simply assume that the reheating temperature is low enough compared with the GUT-scale to make the number of ψ small enough.

2. Model-2

If $m_\psi \ll T_{\text{reheating}}$, large number of ψ can be produced through the reheating process. In the presence of an additional hidden sector fermion field ψ' , we can consider the decay of ψ , $\psi \rightarrow \psi'\phi$. Then, a simple $U(1)_H$ interaction of the form $\bar{\psi}'\psi'\gamma_H$ can induce the pair annihilation of ψ' , $\psi'\psi' \rightarrow \gamma_H\gamma_H$, so that the relic abundance of ψ' could be small enough provided that

$$\langle\sigma v\rangle_{\psi'\psi'\rightarrow\gamma_H\gamma_H} \sim \frac{\pi\alpha_H^2}{m_{\psi'}^2} \gg 2 \times 10^{-9}\text{GeV}^{-2}, \quad (11)$$

for which we assume that the mass of ψ' is relatively light. For instance, if $\alpha_H \sim 10^{-4}$, $m_{\psi'} \ll 4 \text{ GeV}$ insures that the number of relic ψ' is too small to affect the present universe. The least constrained mass range for the hidden sector milli-charged particle is $0.1 \text{ GeV} \lesssim m_{\psi'} \lesssim 500 \text{ GeV}$ [40]. Thus, we can easily find experimentally allowed masses for ψ' satisfying Eq. (11).

IV. CONCLUSION

The recently reported gamma-ray excesses around 130 GeV based on the Fermi-LAT data is difficult to explain with well known dark matter models. We suggest a model with a decaying scalar dark matter ϕ and a heavy hidden fermion ψ charged under a hidden gauge symmetry $U(1)_H$ allowing a Yukawa interaction, $\lambda\phi\bar{\psi}\psi$. In this model, the hidden sector can communicate with the SM sector through the kinetic mixing ($\epsilon F^{\mu\nu}F'_{\mu\nu}$) and the unwanted fast decay to $\gamma\gamma$ is well suppressed by the small mixing ϵ^2 after ensuring the longevity of ϕ by a doubly suppressed loop factor, $\Gamma_\phi/m_\phi \sim (\alpha_H\lambda)^2(m_\phi^2/m_\psi^2)/(256\pi^3)$. Assuming $\alpha_H\lambda \sim 10^{-7}$

and $\epsilon \sim 10^{-2} - 10^{-4}$, we found that the required decay rate $\Gamma(\phi \rightarrow \gamma_H \gamma) \sim 10^{-28} \text{ sec}^{-1}$ for the observed $\sim 130 \text{ GeV}$ gamma-ray flux is obtained with a GUT scale mass for the fermion $m_\psi \sim 10^{16} \text{ GeV}$. The model is free from other observational bounds. We also discussed possible production mechanism for the dark matter from the inflation or a heavy particle decay but other possibilities are still open for studies in the future.

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Appendix A: Kinetic mixing between $U(1)_{\text{EM}}$ and massless $U(1)_H$

Consider two Abelian gauge groups $U(1)_{\text{EM}}$ and $U(1)_H$.⁵ The kinetic mixing between $U(1)_{\text{EM}}$ and $U(1)_H$ is parameterized as

$$\mathcal{L} = -\frac{1}{4} \hat{F}_{\mu\nu} \hat{F}^{\mu\nu} - \frac{1}{4} \hat{F}_{H\mu\nu} \hat{F}_H^{\mu\nu} - \frac{\xi}{2} \hat{F}_{\mu\nu} \hat{F}_H^{\mu\nu}, \quad (\text{A1})$$

where $\hat{A}_\mu (\hat{A}_\mu^H)$ is the $U(1)_{\text{EM}} (U(1)_H)$ gauge boson and its field strength tensor is $\hat{F}^{\mu\nu} (\hat{F}_H^{\mu\nu})$. The kinetic mixing is parameterized by ξ , which is generically allowed by the gauge invariance and the Lorentz symmetry. In the low energy effective theory, the kinetic mixing parameter ξ is considered to be an arbitrary parameter. An ultraviolet theory is expected to generate the kinetic mixing parameter ξ [37].

We are only interested in a small kinetic mixing, $\xi < 1$. Thus, for convenience we can set $\xi \equiv \sin \epsilon$. The kinetic terms for the photon and hidden photon are diagonalized by the following transformation:

$$\begin{pmatrix} A'_\mu \\ A'^H_\mu \end{pmatrix} = \begin{pmatrix} \cos \frac{\epsilon}{2} & \sin \frac{\epsilon}{2} \\ \sin \frac{\epsilon}{2} & \cos \frac{\epsilon}{2} \end{pmatrix} \begin{pmatrix} \hat{A}_\mu \\ \hat{A}_\mu^H \end{pmatrix}. \quad (\text{A2})$$

In this transformed basis, the Lagrangian is given by

$$\mathcal{L} = -\frac{1}{4} F'_{\mu\nu} F'^{\mu\nu} - \frac{1}{4} F'^H_{\mu\nu} F'^{H\mu\nu}, \quad (\text{A3})$$

⁵ For the case of massive $U(1)_H$, see Ref. [39].

where $F'_{\mu\nu}$ and $F'^H_{\mu\nu}$ are the field strength tensors corresponding to A'_μ and A'^H_μ , respectively. Since two gauge bosons are massless, we still have an $SO(2)$ symmetry:

$$\begin{pmatrix} A_\mu \\ A_\mu^H \end{pmatrix} = \begin{pmatrix} \cos \eta & \sin \eta \\ -\sin \eta & \cos \eta \end{pmatrix} \begin{pmatrix} A'_\mu \\ A'^H_\mu \end{pmatrix}. \quad (\text{A4})$$

If we choose the mixing angle as $\eta = -\epsilon/2$, the final relation between the bases (A_μ, A_μ^H) and $(\hat{A}_\mu, \hat{A}_\mu^H)$ is

$$\begin{pmatrix} A_\mu \\ A_\mu^H \end{pmatrix} = \begin{pmatrix} \cos \epsilon & 0 \\ \sin \epsilon & 1 \end{pmatrix} \begin{pmatrix} \hat{A}_\mu \\ \hat{A}_\mu^H \end{pmatrix}. \quad (\text{A5})$$

By this diagonalization procedure of the kinetic terms, we finally obtain

$$\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{4} F_{\mu\nu}^H F^{H\mu\nu}, \quad (\text{A6})$$

where the new field strength tensors are $F_{\mu\nu}$ and $F_{\mu\nu}^H$. The SM photon corresponds to A_μ and the hidden photon to A_μ^H .

Let us take the following simple interaction Lagrangian of a SM fermion f with the photon in the original basis as

$$\mathcal{L}_f = \bar{f} (\hat{e} Q_f \gamma^\mu) f \hat{A}_\mu. \quad (\text{A7})$$

Note that in this basis no direct interaction exists between the SM fermion and the hidden sector gauge boson \hat{A}_H . If there exists a hidden sector fermion ψ with the $U(1)_H$ charge Q_ψ^H , its interaction with the hidden sector gauge boson is simply represented by

$$\mathcal{L}_\psi = \bar{\psi} (\hat{g}_H Q_\psi^H \gamma^\mu) \psi \hat{A}_\mu^H. \quad (\text{A8})$$

In this case, there is also no direct interaction between the hidden fermion and the SM photon \hat{A} .

We can recast the Lagrangian (A7) in the transformed basis (A_μ, A_μ^H) ,

$$\mathcal{L}_f = \bar{f} \left(\frac{\hat{e}}{\cos \epsilon} Q_f \gamma^\mu \right) \psi A_\mu. \quad (\text{A9})$$

Here, one notices that the SM fermion has a coupling only to the visible sector gauge boson A even after changing the basis of the gauge bosons. However, the coupling constant \hat{e} is

modified to $\hat{e}/\cos\epsilon$, and so the physical visible sector coupling e is just defined as $e \equiv \hat{e}/\cos\epsilon$. Similarly, we can derive the following interactions for χ ,

$$\mathcal{L}_\psi = \bar{\psi}\gamma^\mu (\hat{g}_H Q_\psi^H A_\mu^H - \hat{g}_H \tan\epsilon Q_\psi A_\mu) \psi. \quad (\text{A10})$$

In this basis, the hidden sector matter field ψ now can couple to the SM photon A with the coupling $-\hat{g}_H Q_\psi^H \tan\epsilon$. Consequently, we can interpret the hidden particle ψ as a particle with a EM charge $Q_\psi \equiv (-\hat{g}_H Q_\psi^H \tan\epsilon)/e$. In addition, we can set the physical hidden sector coupling g_H as $g_H \equiv \hat{g}_H$.

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